

The Role of Partial Ionization Effects in the Chromosphere

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The energy for the coronal heating must be provided from the convection zone. However, the amount and the method by which this energy is transferred into the corona depends on the properties of the lower atmosphere and the corona itself. We review: 1) how the energy could be built in the lower solar atmosphere; 2) how this energy is transferred through the solar atmosphere; and 3) how the energy is finally dissipated in the chromosphere and/or corona. Any mechanism of energy transport has to deal with the various physical processes in the lower atmosphere. We will focus on a physical process that seems to be highly important in the chromosphere and not deeply studied until recently: the ion-neutral interaction effects in the chromosphere. We review the relevance and the role of the partial ionization in the chromosphere and show that this process actually impacts considerably the outer solar atmosphere. We include analysis of our 2.5D radiative MHD simulations with the *Bifrost* code (Gudiksen *et al.* 2011) including the partial ionization effects on the chromosphere and corona and thermal conduction along magnetic field lines. The photosphere, chromosphere and transition region are partially ionized and the interaction between ionized particles and neutral particles has important consequences on the magneto-thermodynamics of these layers. The partial ionization effects are treated using generalized Ohm's law, i.e., we consider the Hall term and the ambipolar diffusion (Pedersen dissipation) in the induction equation. The interaction between the different species affects the modeled atmosphere as follows: 1) the ambipolar diffusion dissipates magnetic energy and increases the minimum temperature in the chromosphere; 2) the upper chromosphere may get heated and expanded over a greater range of heights. These processes reveal appreciable differences between the modeled atmospheres of simulations with and without ion-neutral interaction effects.

Keywords: Magnetohydrodynamics MHD — **Methods:** numerical — **Sun:** atmosphere — **Sun:** magnetic field

1. Coronal heating vs heating concentrated in the low atmosphere

One of the recent and heated discussions on chromospheric and corona heating concerns the location of magnetic energy release in the solar atmosphere. Is it

highly dependent on height? What is the role of flux emergence and what is the role of the magnetic field topology?

One dimensional hydrodynamic simulations can cast some light on this issue. For example, Testa *et al.* (2014) use 1D models to prove that nano-flares situated in the corona can spread energy via non-thermal electron beams rather than through thermal conduction. As another example, Klimchuk & Bradshaw (2014) compared synthetic data obtained from 1D multi parametric hydro simulations with the Atmospheric Imaging Assembly observations (AIA, Lemen *et al.* 2012; Boerner *et al.* 2012) on board of the Solar Dynamics Observatory (SDO, Pesnell *et al.* 2012). They claimed that their synthetic data can not reproduce the observables unless they consider quasi-steady coronal heating or coronal nano flares situated well above the chromosphere. This particular case highlights the limits of a 1D hydrodynamic approach: advanced 3D simulations of heating in association with chromospheric jets indicate that the complex 3D topology needs to be taken into account to understand the impact on the corona (Hansteen *et al.* 2010).

It is clear that studying the emergence and topology of the magnetic field and the likely sites of energy release requires full 2D or 3D models covering the region from the photosphere to the corona. Sophisticated simulations that include thermal conduction along magnetic field lines (Gudiksen & Nordlund 2004), radiative losses in the corona and chromosphere (Hansteen *et al.* 2007) are now well established tools and provide quite sophisticated comparisons with observations (Hansteen *et al.* 2010; Bingert & Peter 2011). These models show that the average heating per particle coming from the released magnetic energy is largely concentrated in the region between the upper chromosphere and the lower corona. Synthetic observations generated from these models reveal overall similarities with observed intensities, line shifts, widths, and profiles (see e.g., Gudiksen & Nordlund 2005*b*; Peter *et al.* 2006; Olluri *et al.* 2013*b*; Hansteen *et al.* 2014). For instance, Hansteen *et al.* (2010) use the Bifrost code (Gudiksen *et al.* 2011) and achieve rather good agreement between synthetic and observed average redshifts for transition region spectral lines. They argue that this is due to episodic heating being localized in the lower atmosphere (somewhere between the middle chromosphere and the lower corona).

Simulations by Martínez-Sykora *et al.* (2011*a*) also reveal another observable due to the presence of localized episodic heating: they studied the asymmetries of the EUV spectral profiles using so-called Red-Blue asymmetry (RB) analysis (De Pontieu *et al.* 2009). The largest blue-ward asymmetries are for transition region EUV profiles (Kjeldseth Moe & Nicolas 1977; Peter 2001; Hara *et al.* 2008; De Pontieu *et al.* 2009; McIntosh & De Pontieu 2009; Peter 2010; Tian *et al.* 2011). Martínez-Sykora *et al.* (2011*a*) find that the synthetic RB analysis of various EUV lines from their simulations reproduces, qualitatively, the observations taken by Solar Ultraviolet Measurements of Emitted Radiation (SUMER, Wilhelm *et al.* 1995) on board of the Solar and Heliospheric Observatory (SOHO, Domingo *et al.* 1995) and the EUV Imaging Spectrometer (EIS, Culhane *et al.* 2000) on board of Hinode (Kosugi *et al.* 2007). Furthermore, they find that the episodic heating produces a larger range of velocities within the transition region than in the corona. Moreover, the simulations by Martínez-Sykora *et al.* (2011*a*) include small scale flux emergence, and find that the average volumetric heating is still localized largely within the transition region (Martínez-Sykora *et al.* 2009*a*; Hansteen *et al.* 2010). Despite this, they find a quantitative disagreement with observations when looking at line

asymmetries, most likely due to the lack of vigorous dynamics in the simulations. A word of caution is in order, these synthetic spectral results and models may be significantly impacted by the simplified magnetic field configuration utilized in the simulations. The penalty paid in performing full 3D simulations is the high resistivity required to keep structures larger than the numerical grid scale. Small scale plasma physics can therefore not be fully resolved, and the treatment of phenomena such as the damping of shocks and the magnetic reconnection process occur at scales much larger than in nature. Therefore, one must be aware of the limitations of not resolving some of the physical processes which may impact the transport of the magnetic field stresses (built by the convective motions) to greater heights in the upper atmosphere.

2. How to build stress in the atmosphere

Let us start at the source: it is well known and broadly accepted that convective motions build stresses in the magnetic field in the atmosphere. In the quiet sun and coronal holes, there are at least five different mechanisms in the photosphere that could build stress or transport energy into the upper layers. The first is the magneto-acoustic shocks produced by the convective motions, p-modes, overshooting, convective collapse, etc (Carlsson & Stein 1992, 1995; Bellot Rubio *et al.* 2001; Stein & Nordlund 2006; Skartlien *et al.* 2000; Hansteen *et al.* 2006; De Pontieu *et al.* 2007*a*; Martínez-Sykora *et al.* 2009*b*; Kato *et al.* 2011 among many others), or even by reconnection or magnetic energy release in the photosphere or lower chromosphere processes (Martínez-Sykora *et al.* 2009*b*). The second lies in waves, where the restoring force is “magnetic” such as kink-, Alfvén-, or other transversal waves, in the following called “alfvenic”, produced by reconnection processes and/or convective motions (De Pontieu *et al.* 2007*b*; Jess *et al.* 2009; Goossens *et al.* 2009; McIntosh *et al.* 2011; van Ballegoijen *et al.* 2011; De Moortel & Pascoe 2012; De Pontieu *et al.* 2014). The third is by small scale flux emergence (Centeno *et al.* 2007; Martínez González & Bellot Rubio 2009; Martínez González *et al.* 2012; Sainz Dalda *et al.* 2012 among others). In short, these observations reveal that small scale loops occur throughout the photosphere, e.g., Martínez González *et al.* (2012) observations of the quiet sun revealed the emergence of small scale magnetic field loops in the photosphere at a rate of the order of $0.2 \text{ loops h}^{-1} \text{ arcsec}^{-2}$. The fourth mechanism is through mass flux. The proponents note that mass flux also can contribute to transport or build energy in the upper layers of the atmosphere, for instance photospheric supersonic jets observed by Bellot Rubio (2009) with the Solar Optical Telescope (SOT, Tsuneta *et al.* 2008) on board of Hinode and by Borrero *et al.* (2012) with IMAX on board of the balloon-borne solar observatory SUNRISE (Solanki *et al.* 2010). Another example of mass flux contributing to energy supply in the outer layers is that of spicules, with De Pontieu *et al.* (2011) suggesting that some of the spicular mass flux is heated to coronal temperatures, based on observations with Hinode/SOT and SDO/AIA (see also Martínez-Sykora *et al.* 2011*b*, 2013). Finally, the motion of magnetic flux elements in the photosphere, driven by convective motions will easily build magnetic field stresses that propagate into the corona, as illustrated with the well known cartoon by Parker (1983). Note, these mechanisms are not completely independent, e.g., flux emergence can also drive waves. Another

possible source of the mechanical energy needed to build extra current is the highly dynamic chromosphere (see below).

We would like to describe in greater detail two of the sources listed above that we consider of great interest in this review. The first one concerns small scale flux emergence. Small bipole structures in the Stokes V magnetograms suggest tiny loops within the granules such as observed by Lites *et al.* (1996, 1998), De Pontieu (2002) and later by Centeno *et al.* (2007), Ishikawa *et al.* (2008), Martínez González & Bellot Rubio (2009), and Gömöry *et al.* (2010), among others. A large number of small scale loops are found to occur in the quiet sun. In addition, strong asymmetric Stokes V profiles suggest structures of unresolved, even tinier, loops (Grossmann-Doerth *et al.* 2000; Sainz Dalda *et al.* 2012; Viticchié 2012). Sainz Dalda *et al.* (2012) combined inversion with synthetic profiles from 3D radiative MHD simulations in order to explain the strong asymmetry profiles. Viticchié *et al.* (2011) categorized a large sample of different asymmetries and inverted the corresponding Stokes profiles. Many of these strong asymmetries suggest the presence of two different polarities of magnetic field within the same pixel. Finally, various types of observations reveal pervasive large amount of horizontal magnetic flux at the photosphere (Lites *et al.* 1996, 2008; Borrero & Kobel 2011; Bellot Rubio & Orozco Suárez 2012; Stenflo 2013 among others). Martínez González *et al.* (2012) suggested that 30% of these tiny loops are able to reach the chromosphere (observations, for instance, by Centeno *et al.* 2007; Martínez González & Bellot Rubio 2009; Guglielmino *et al.* 2010; Ortiz *et al.* 2014 also reveal the various chromospheric signals of flux emergence).

The second aspect we would like to discuss in a bit more detail is the motion of the magnetic flux elements. Magnetic field elements move around due to the convective motions, collapsing granules and emerging new convective cells. The magnetic field elements evolve with the convective motions, merging elements of the same polarity, fragmenting magnetic elements, and canceling magnetic elements of different polarity (Schrijver *et al.* 1997). It is important to understand how the convective motion is braiding and tangling the magnetic field in the solar atmosphere (Lamb *et al.* 2008; DeForest *et al.* 2007; Zhou *et al.* 2010; Giannattasio *et al.* 2014 among others) since the convective motion could build magnetic energy and thereafter this energy may dissipate into kinetic and thermal energy. These authors studied the dynamic properties of the magnetic field elements at the photosphere in the quiet sun. Most of the identified magnetic network elements seem to form due to unresolved small scale shredding, coalescence of previous magnetic flux, magnetic field cancelation and emergence.

Moreover, it is crucial to calculate the magnetic elements that have merged, canceled and/or fragmented in order to understand how the magnetic field is stressed in the atmosphere (Schrijver *et al.* 1997; Iida *et al.* 2012; Zhou *et al.* 2013; Gošić *et al.* 2014). Gošić *et al.* (2014) using long time series of Hinode/SOT observations followed and identified different internetwork elements, and noting which of these reach the network. Their conclusions are that internetwork elements interact with network elements and thereby modify its net magnetic flux. Only a tiny part of the internetwork elements that interact with the network comes from flux emergence and most of the network flux comes from shredding and coalescence of previous internetwork magnetic flux. The internetwork magnetic flux replaces the network flux in a time period of 9 hours. On average, 40% of the internetwork individual elements ends up, eventually, in the network (Gošić *et al.* 2014).

3. The various “filters” of the various atmospheric layers

As discussed above, convective motions have several methods of generating enough energy to heat the corona. The next question becomes: What happens to this energy for each source mentioned in the previous section as it propagates through the solar atmosphere? In the solar atmosphere there are many physical processes that play different roles in different layers. Likewise, one must also consider whether the transport of energy can trigger other mechanisms that contribute further to increasing the magnetic stresses in the atmosphere. Finally, there remains the question of where magnetic energy is released for each source, if it is released at all.

In short, in order to reach the corona, the energy must go through:

- The photosphere, where a large amount of the convective energy is released by radiative transport, and the atmosphere is highly stratified with a pressure scale height of the order of less than 100 km. The upper photosphere is dominated by large velocities due to convective overshoot.
- The chromosphere is the interface layer between the solar surface (photosphere) and the million degree corona. The chromosphere is of great interest because, in principle, it contains enough non-thermal energy to heat the entire corona. Many complex physical processes play a role in the chromosphere: *a)* the plasma is in non-local thermodynamic equilibrium (NLTE), *b)* the radiation is optically thick, *c)* radiation is highly scattered, *d)* ionization is not necessarily in equilibrium, *e)* the gas is partially ionized, *f)* in the upper chromosphere, transition region and corona, the thermal conductive flux propagates along the magnetic field lines. In addition, the chromosphere is also distinguished by several transitions such as from non-magnetized to magnetized, and from partial to full ionization, from high to low plasma β (where plasma $\beta = \frac{4\pi P}{B^2}$ is the ratio between the gas pressure and magnetic pressure), from optically thick to optically thin, etc. Note that many of the processes (especially from the items *a* to *d*) lead to highly complex interpretations of imaging and spectral observations (e.g., Leenaarts *et al.* 2012, 2013).
- The transition region, located between the chromosphere and corona, is optically thin, magnetically dominated, with steep temperature gradients, filled with shocks and flows approaching or exceeding the local speed of sound, and where, as mentioned, thermal conduction is a highly important process. In this region we cannot expect ions to be in ionization equilibrium and for low density regions such as coronal holes we may also find that ion and electron temperatures differ (Hansteen *et al.* 1993; Lie-Svendsen *et al.* 2001).
- Finally, in the corona, which generally has low plasma beta, radiation is optically thin, and, as in the transition region, the ionization of some heavy ions is not in equilibrium (Joselyn *et al.* 1977; Hansteen *et al.* 1993; Bradshaw & Cargill 2006; Olluri *et al.* 2013*a*), background radiation heats lower layers, the timescales of several important processes are very short and non-thermal physics may be vital in describing these processes.

Any kind of magnetic stress built by the convective motions that reaches the corona will be strongly affected by the various processes mentioned above. For instance, the chromosphere appears to be filled with magnetic field, even in the quiet sun, despite the fact that the photosphere is sub-adiabatic, which thwarts the expansion of emerging magnetic flux into the regions above (Acheson 1979; Archontis *et al.* 2004). Therefore, it remains unknown under which conditions small-scale flux emerges into the chromosphere. The challenge facing the researcher is thus to combine all these ingredients into a model of the entire solar atmosphere, taking into account all important effects while discarding those that have little impact on the buildup, transport and dissipation of the energy flux heating the outer atmosphere.

4. Partial ionization effects

In this paper we will concentrate on ion-neutral interaction effects among all the possible physical processes from the photosphere to the lower corona. The chromosphere is partially ionized: ions are coupled to the magnetic field, whereas neutrals are not directly affected by the presence of the magnetic field and they can move “freely”. Sufficient collisions between ions and neutrals will couple the neutrals to the magnetic field, while at the same time to some extent allowing ions to slip across the field. As a result, the magnetic field can diffuse and magnetic energy will be dissipated into thermal energy. In order to model this scenario, the fluid should be treated as consisting of two or three fluids, i.e., ion, neutral and/or electron particles. To solve the equations for all three species is computationally highly demanding. Therefore two or three fluid codes focus on specific and localized problems (Smith & Sakai 2008; Sakai & Smith 2009; Meier 2011; Leake *et al.* 2012 among others) instead of modeling large portions of the solar atmosphere, i.e., instead of including the convection zone, photosphere, chromosphere, transition region and corona. Additionally these codes rarely consider many of the other important physical processes in the solar atmosphere such as radiative transfer, thermal conduction along the magnetic field lines, etc, which is the matter of interest we wish to cover in this review (see previous sections).

(a) *The Hall Term and Ambipolar Diffusion*

Fortunately, the chromosphere seems to be sufficiently collisional to allow us to convert the three fluid problem into a single MHD fluid while still taking into account the most important partial ionization effects. In order to do this one assumes that 1) collisions are sufficiently numerous and the ion and electron temperatures are the same, 2) the relevant resolvable timescales one wants are assumed to be greater than the collision times between ions and neutrals (Cowling 1957; Braginskii 1965; Parker 2007; Pandey & Wardle 2008; Leake *et al.* 2013a). According to the comparison of the collision rates with timescales of various relevant physical processes from the self-consistent 2D radiative-MHD atmospheric models done by Martínez-Sykora *et al.* (2012) using the Bifrost code (Gudiksen *et al.* 2011), the approximations mentioned above seem to be fulfilled in the chromosphere most of the time.

In this case, the induction equation expands into the so-called generalized induction equation which includes at least two new terms, i.e., the Hall term and the ambipolar diffusion (or Pedersen dissipation) term (Cowling, 1957; Braginskii, 1965) as follows:

$$\frac{\partial \mathbf{B}}{\partial t} = \nabla \times \left[\mathbf{u} \times \mathbf{B} - \eta_{ohm} \mathbf{J} - \frac{\eta_{hall}}{|\mathbf{B}|} \mathbf{J} \times \mathbf{B} + \frac{\eta_{amb}}{B^2} (\mathbf{J} \times \mathbf{B}) \times \mathbf{B} \right] \quad (4.1)$$

where \mathbf{B} , \mathbf{J} , \mathbf{u} , and η_{ohm} are magnetic field, current density, velocity field, and the ohmic diffusion, respectively. The new “diffusion” terms are:

$$\eta_{ohm} = \frac{m_e(\nu_{ei} + \nu_{en})}{q_e^2 n_e} \quad (4.2)$$

$$\eta_{hall} = \frac{|\mathbf{B}|}{q_e n_e} \quad (4.3)$$

$$\eta_{amb} = \frac{(|\mathbf{B}| \rho_n / \rho)^2}{\rho_i \nu_{in}} = \frac{(|\mathbf{B}| \rho_n / \rho)^2}{\rho_n \nu_{ni}} \quad (4.4)$$

where ρ_i , ρ_n , ρ , ν_{in} , ν_{ni} , n_e , and q_e are ion density, neutral density, total density, ion-neutral collision frequency, neutral-ion collision frequency, electron number density and electron charge, respectively. Note that η_{hall} is not a diffusive or dissipative term in contrast to η_{ohm} and η_{amb} . The ambipolar diffusion should not be confused with the so-called ambipolar drift used in plasma physics. For plasma physics, ambipolar drift or electric field is due to the relative velocity or temperature difference between ions and electrons, e.g., Arefiev & Breizman (2008). Ambipolar diffusion is also referred to as Pedersen dissipation.

In short, the ohmic diffusion can change directly the magnetic field topology, i.e., change the connectivity of the field lines with reconnection, but the Hall and the ambipolar terms cannot, without the help of ohmic diffusion, since the former term does not follow the formalism of $\mathbf{u} \times \mathbf{B}$ (Biskamp 2005) and the other two follow it, where \mathbf{u} for the Hall term is \mathbf{J} and for the ambipolar term is $\mathbf{J} \times \mathbf{B}$. The ohmic (first term in the right hand side of the following expression) and ambipolar diffusion (second term in the right hand side of the following expression) can dissipate magnetic energy into thermal energy as follows:

$$\frac{\partial e}{\partial t} \propto \eta_{ohm} J^2 + \eta_{amb} \frac{|\mathbf{J} \times \mathbf{B}|^2}{|\mathbf{B}|^2} \quad (4.5)$$

(see Parker 2007 for details).

The physical process behind the Hall term is that in a fully ionized plasma as in the corona, the relative immobility of ions relative to electrons in response to the Lorentz force leads to an electric field parallel to $\mathbf{J} \times \mathbf{B}$. The Hall term for weakly ionized plasma is created due to collisions between ions and neutrals. The neutrals decouple ions from the magnetic field lines leading again to an electric field parallel to $\mathbf{J} \times \mathbf{B}$. The physical process behind the ambipolar diffusion is that neutrals do not experience the Lorentz force. As a result, neutrals decouple from the magnetic field and allow the magnetic field to diffuse through the neutral gas (Braginskii 1965).

(b) *Impact of the Ion-Neutral Interaction Effects on Simplified Physical Processes*

Partial ionization effects impact various physical processes that play a role in the chromosphere and prominences. For instance, the partial ionization effects in the chromosphere are known to allow ion-neutral collisions to dissipate alfvénic waves (De Pontieu & Haerendel 1998; De Pontieu 1999). However, the picture is far from complete, since it is unknown if this dissipation plays a major role in heating the chromosphere (Biermann 1948; Leake *et al.* 2005; Hasan & van Ballegoijen 2008; Song & Vasyliūnas 2011; Madsen *et al.* 2014). Most of these studies obtain different high frequency cut-offs for alfvénic waves for different features (spicules, filaments, prominences, thin flux tubes, etc) using fixed values for the ion-neutral collision frequency of semi-empirical models such as the VAL-C model (De Pontieu *et al.* 2001; Leake *et al.* 2005; Forteza *et al.* 2007). This dissipation contributes to the thermal and dynamic energy budget. In fact, as mentioned in the previous section, electrical currents perpendicular to the magnetic field can be dissipated by ambipolar diffusion and convert magnetic energy into thermal energy (Arber *et al.* 2009; Goodman 2011; Khomenko & Collados 2012; Goodman & Judge 2012). Soler *et al.* (2009, 2012) deduced that kink waves are damped by diffusion parallel to the magnetic field and Alfvén waves by perpendicular diffusion. For short wavelengths, kink waves are converted into Alfvén waves due to the collisions between ion and neutrals. Ohmic dissipation dominates for long wavelengths. In fact, all the diffusive mechanisms must be considered, and missing any of the inertial, Hall and ambipolar terms will give incorrect results (Khodachenko *et al.* 2004, 2006). Zaqarashvili *et al.* 2012 deduced that there is no cut-off frequency for Alfvén waves, only a damping mechanism due to the ion-neutral interaction effects, when one is taking into account both the inertial and Hall term (where the inertial term is defined as the time variation of the relative velocity perturbations perpendicular to the unperturbed magnetic field between ions and neutrals). However, according to Vranjes & Kono (2014), this result might be due to the assumption of neglecting the exchange of momentum between ions and neutrals. However, one must be really careful with all these results since the ion-neutral collision frequency has been calculated in a very simplified manner and a proper treatment of the collision between ion and neutrals may provide different results (see below). In fact, even in semi-empirical models, using different methods to calculate the ion-neutral collision frequency provide very different values for the damping and ambipolar diffusion terms (De Pontieu *et al.* 2001).

Ambipolar diffusion allows magnetic field to diffuse into the atmosphere (Leake & Arber 2006; Arber *et al.* 2007; Leake & Linton 2013). Since neutrals move through the expanding magnetic field in the chromosphere, the magnetic field lines lift less matter. As a result the ambipolar diffusion inhibits the Rayleigh-Taylor (RT) instability. In contrast, a completely different configuration, such as the 2.5D simulations of prominences including partial ionization, shows an increase of small scale velocities as a result of the non-linearity of the RT instability, in contrast to single fluid simulations (Díaz *et al.* 2014; Khomenko *et al.* 2014). Note that, the ambipolar diffusion can have effects of opposite nature: it can either inhibit the RT instability or increase the small scale structures of the RT instability depending on the spatial distribution of ambipolar diffusion, i.e., compare Arber *et al.* (2007) with Díaz *et al.* (2014) and Khomenko *et al.* (2014). Since the ambipolar diffusion is highly depen-

dent on the thermal properties of the plasma, the ion-neutral collision frequency needs to be calculated directly from the thermal properties of the plasma and the energy balance must be solved self-consistently.

Regardless of the assumptions of the background atmosphere and the calculation of the ion-neutral collision frequency, it is clear that partial ionization effects play an important role in these processes.

As mentioned before, ion-neutral interaction effects cannot change the topology of the magnetic field, but they can impact the reconnection rate. On one hand, ambipolar diffusion causes sharp structures to evolve around the reconnection region, thus setting the scene for tearing reconnection, accelerating the reconnection rate or, depending on the configuration of the magnetic field and the diffusivities, singularities that initiate reconnection. On the other hand, electron pressure and ohmic dissipation act against the formation of sharp currents, i.e., the reconnection rate depends on the ionization rate and ion-neutral collision frequency (Brandenburg & Zweibel 1994, 1995). When the ionization decreases, the reconnection rate decreases too (Smith & Sakai 2008; Sakai & Smith 2009). Note that the two fluid simulations from the latter two references used fixed ionization fractions and did not take into account the ionization/recombination rates. In addition, they did not clarify if their parametric study is for a fixed density or a fixed number of ions since an increase of density can lead to a decrease of the reconnection rate. In a reconnection process in a weakly ionized plasma when ions and neutrals are decoupled, it is found that an excess of ions can accumulate in the reconnection region, as shown with the simulations done by Vishniac & Lazarian (1999) and Lazarian *et al.* (2004). The recombination in the reconnection region removes ions which combined with the alfvénic flows can lead to fast reconnection. This has been tested with 2D two fluid simulations taking into account the ionization/recombination, optically thin radiative transfer and collisional heating by Leake *et al.* (2012, 2013*b*). An important remark from the latter authors is that it is crucial to treat the ambipolar diffusion while at the same time solving the energy balance including the ionization/recombination since it can change drastically the impact of the ambipolar diffusion on the reconnection process. While the results described above indicate that ion-neutral effects are an important ingredient in a variety of processes in the solar atmosphere, the current simplified approaches to the energy balance significantly limit the applicability of these results in the solar atmosphere.

All the studies mentioned in this section are focussed on specific driving processes. These studies are extremely useful to discern the physical mechanism in each specific problem. However, in order to place these processes in the context of the solar atmosphere, we need to combine this knowledge with self-consistent radiative MHD models that have the proper convective drivers (see Section 2) and that include the relevant physical processes in each layer of the atmosphere, as done by Hansteen *et al.* (2007), Abbett (2007), Martínez-Sykora *et al.* (2008), and Hansteen *et al.* (2010), amongst others. The missing ingredient in these models has so far been ion-neutral interactions which should be calculated self-consistently with the thermal properties (such as done by Leake *et al.* 2012, 2013*b* for simplified magnetic reconnection simulations). The Hall effect in the photosphere has been implemented following this approach by Cheung & Cameron (2012). The preliminary work done in Martínez-Sykora *et al.* (2012) added ambipolar diffusion to the

description of the outer solar atmosphere. We describe this type of simulations including ambipolar diffusion in the next sections.

(c) *Importance of the Ambipolar Diffusion in the Chromosphere*

In the previous section we described several aspects of how ion-neutral interaction impacts various physical processes in the chromosphere and prominences. Let us now discuss the range of values for the various “diffusion” terms of the expanded induction equation in the lower solar atmosphere. Figure 1 in Khomenko & Collados (2012) shows, using the semi-empirical VAL-C model and thin magnetic flux tube approximation (Spruit 1981; Roberts & Webb 1978), the values of the various “diffusion” terms through the atmosphere. In short, ohmic diffusion dominates below the photosphere. In the photosphere, the Hall term is the largest and reaches $7 \times 10^5 \text{ m}^2 \text{ s}^{-1}$. In the chromosphere the largest, by roughly 2 orders of magnitude, is the ambipolar diffusion ($10^8 \text{ m}^2 \text{ s}^{-1}$). However, these values become extremely different if we consider a radiative MHD simulation such as the one used by, for instance, Gudiksen & Nordlund (2005a), Hansteen *et al.* (2010) (see below).

In order to calculate the ambipolar diffusion terms, the ion-neutral collision frequency must be obtained from the ionization state and thermal properties of the plasma. For this there are several methods (Osterbrock 1961; von Steiger & Geiss 1989; Fontenla *et al.* 1993 among others). Osterbrock (1961) uses elastic scattering cross sections, von Steiger & Geiss (1989) considers hard spheres, and Fontenla *et al.* (1993) uses different sets of elastic collision cross sections than Osterbrock (1961) to determine the ion-neutral collision frequency. The final results of these methods differ in the ion-neutral collision frequency by an order of magnitude for semi-empirical models (De Pontieu *et al.* (2001) among others). In a more realistic approach, Martínez-Sykora *et al.* (2012) show that the various methods used to calculate the ion-neutral collision frequency differ most in chromospheric regions where the ambipolar diffusion is largest. A more recent analysis of the ion-neutral collision frequency is done by Vranjes & Krstic (2013) where they include accurate cross sections which vary as a function of temperature, due to quantum effects, and between descriptions of elastic scattering, momentum transfer and viscosity. As mentioned by these authors, this work can be improved including the effects of inelastic collisions.

Figure 1 shows the ohmic diffusion, Hall term and ambipolar diffusion where Osterbrock (1961) was used to calculate the ion-neutral collision frequency. The simulation is 2.5D and ranges from the upper convection zone to the lower corona including radiative transfer with scattering from the photosphere to the corona (Hayek *et al.* 2010; Carlsson & Leenaarts 2012), thermal conduction along the magnetic field lines and partial ionization effects using the Bifrost code (see Martínez-Sykora *et al.* 2012 and references cited within). The initial magnetic field is unipolar with a mean unsigned field strength of 5 G in the corona. The convective motions build enough magnetic field stress to self-consistently maintain the hot corona (Galsgaard & Nordlund 1996; Gudiksen & Nordlund 2004, 2005b; Hansteen *et al.* 2010; Martínez-Sykora *et al.* 2011a). Note that the magnetic field is vertical in the corona (left panel), and despite the fact that our simulations do not include imposed flux emergence, the dynamics in the simulation and the convective motions accumulate horizontal magnetic field in the sub-adiabatic photosphere. As mentioned above, this horizontal field

will in principle be constrained to the sub-adiabatic photosphere (Acheson 1979; Archontis *et al.* 2004). In addition, magnetic field footpoints will move around. Note that these simulations naturally reproduce several of the energy generating drivers discussed in Section 2. Moreover, the energy flux through the atmosphere will be modified by the physical processes considered in the model as it propagates upwards towards the corona as mentioned above (see Martínez-Sykora *et al.* 2012 and compare with Section 3).

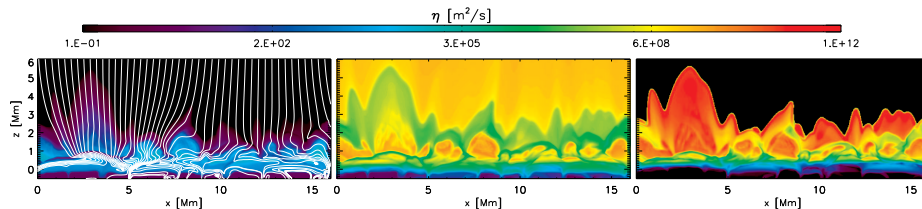


Figure 1. Maps of ohmic diffusion (left), Hall term (middle) and ambipolar diffusion (right) from a snapshot of a 2D radiative-MHD simulation that includes partial ionization effects. The color scheme is on a logarithmic scale. The magnetic field is drawn with white lines in the left panel.

The ambipolar diffusion is the largest term in the induction equation in the chromosphere by almost ten orders of magnitude compared to the ohmic dissipation and something between two and six orders of magnitude larger than the Hall term (see Figure 1 and Martínez-Sykora *et al.* 2012, the latter compares these terms with the hyper-diffusion intrinsic in the code too.). In addition, the various terms show large spatial variability. For instance, the ambipolar diffusion changes at very small scales (within a few hundred km) by up to seven orders of magnitude, e.g., at height $z = 1$ Mm. This large variability is due to the large changes in the ion-neutral collision frequency and in the neutral density when the ionization state and thermal properties of the fluid are taken into account. Note that this model is self-consistent and the ionization state and thermal properties are set by the various drivers, such as shocks, and convective motions of the footpoints that build magnetic field stress (Section 2) and the physical processes implemented in the simulation: radiation, thermal conduction along the magnetic field lines, scattering, partial ionization effects in the equation of state and ion-neutral interaction effects. One important process that has yet to be taken into account in these models which will change the ohmic diffusion, Hall term, and ambipolar diffusion is the effect of time dependent non-equilibrium ionization of hydrogen and helium (Leenaarts *et al.* 2007; Golding *et al.* 2014) which are known to be important in the upper chromosphere. In our models we assume time dependent equilibrium ionization. These effects must be included in order to draw firm conclusions on the importance of ion-neutral interactions in this part of the solar atmosphere.

From Figure 1, it is clear that any magnetic energy flux (driven by photospheric or chromospheric processes) that is transferred through the chromosphere may be strongly influenced and/or diffused by the ambipolar diffusion. We can therefore already conclude that the chromosphere may diffuse more magnetic field stress than any other atmospheric layer. Note that the ohmic diffusion is greater in the photosphere and chromosphere than in the corona.

Table 1. *Description of the simulations*

(The left column lists the names of the various 2D simulations, and the right column gives a short description of each simulation.)

Name	Description
NGOL	without ion-neutral interaction effects
GOL-OS	with ion-neutral interaction effects using Osterbrock (1961) for ν_{in}
GOL-F	with ion-neutral interaction effects using Fontenla <i>et al.</i> (1993) for ν_{in}

(d) *Impact on the Solar Atmosphere*

Here we will not describe in detail how ambipolar diffusion affects the various processes such as flux emergence and wave propagation in 2D radiative MHD simulations. This is outside of the scope of this review and will be described in a series of follow up papers. However, we will show the impact of ambipolar diffusion on the thermal properties of the solar atmosphere. For this we compare different 2D radiative-MHD simulations listed in table 1.

The different models, i.e., NGOL, GOL-OS and GOL-F, show important differences. This means that even with the large artificial diffusion intrinsic in these codes, the resulting models can discern (at least) some of the effects produced by the partial ionization interaction effects. A first look at the 2D temperature maps in Figure 2 reveals at least two thermal properties that differ between the various simulations: 1) the cold chromospheric expanding bubbles have higher temperatures in the GOL simulations, and they do not rise as high as in simulations without partial ionization effects. 2) The transition region shows a smaller temperature gradient (i.e., is spread out over a larger height range) in the GOL-OS simulation than in the other two simulations.

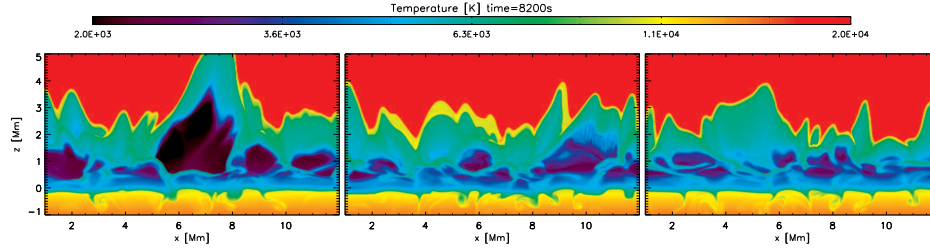


Figure 2. Temperature maps for the NGOL simulation (left), GOL-OS (middle) and GOL-F (right) reveal differences in the thermal properties. The temperature is shown in logarithmic scale.

Cold chromospheric bubbles are formed by expanding shock fronts that pass through the chromosphere. These bubbles contain the lowest temperatures in the atmospheric model (Leenaarts *et al.* 2011). Statistically, the simulations with partial ionization effects have hotter and denser cold chromospheric expanding bubbles than the NGOL simulation. This is shown in the Joint Probability Distribution Functions (JPDF) of the density and temperature integrated over 30 minutes in Figure 3 for simulations NGOL (left panel), GOL-OS (middle panel), and GOL-F (right panel). Everything below 4000 K ($\log T < 3.6$) corresponds to the cold chromospheric bubbles. The minimum temperature achieved is lower in the NGOL

simulation than in the GOL simulations (enhanced with white circle E in the middle panel). In fact, the coolest temperatures for the NGOL simulation are controlled by an *ad-hoc* heating term in the Bifrost code which is introduced to avoid reaching temperatures lower than ~ 1600 K since below this threshold the tabulated equation of state is no longer accurate (Carlsson & Leenaarts 2012). While this *ad-hoc* heating is crucial for the NGOL simulation, it is rarely necessary in both simulations with ion-neutral interaction effects since the cold chromospheric bubbles are not cold enough to reach the temperature at which the *ad-hoc* heating starts. The cold chromospheric bubbles in the simulations with partial ionization effects are a few hundred of degrees hotter than in the NGOL simulation. Leenaarts *et al.* (2011) found no mechanism to prevent extremely low temperatures from occurring in expanding bubbles within the standard, single-fluid MHD formulation. The joule heating contribution coming from ambipolar diffusion provides such a mechanism. Observations can help determine the amount and temperature of such cold plasma in the solar chromosphere (Ayres & Rabin 1996; Penn *et al.* 2011). The details of the physical processes in the chromospheric cold bubbles will be described in a follow up paper.

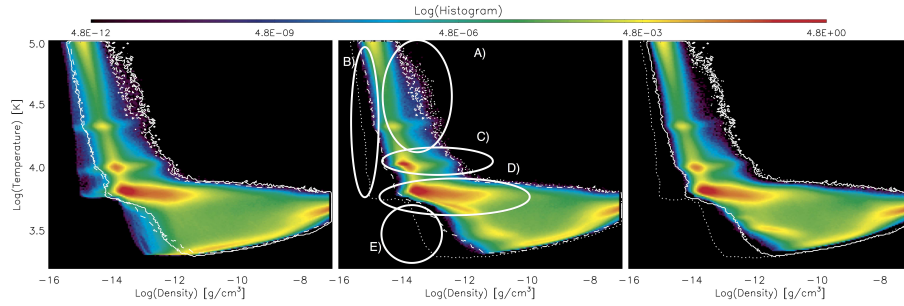


Figure 3. Joint Probability Density Function (JPDF) of temperature (vertical axis) and density (horizontal axis) integrated over 30 min for simulations NGOL (left), and GOL-OS (middle) and GOL-F (right). The white contours correspond to the temperature and density regime of the simulation NGOL in dotted line, GOL-OS in solid line and GOL-F in dashed line in order to simplify the comparison between the simulations. The circles in the middle panel enhance the differences between the various simulations.

The upper chromosphere also shows considerable differences between the three simulations. One can appreciate that in the GOL simulations there are a greater number (compared to NGOL) of grid points concentrated around the temperature where hydrogen ionizes in the simulations (white circle D in the middle panel of Figure 3). Moreover, the number of grid points around the temperature where helium ionizes ($T \sim 10^4$ K) is greater for GOL-OS compared to the other simulations (circle C). In fact, GOL-OS shows several grid points with rather dense material in the transition region (circle A) while the transition region is the least dense in NGOL (circle B). GOL-F has the narrowest range of densities in the transition region. These discrepancies in the GOL simulations are due to the differences in the distribution of ambipolar diffusion as a result of the different methods to calculate the ion-neutral collision frequency and cross section values. hydrogen and helium time dependent ionization may spread the amount of plasma over a larger range

of temperatures than for the case of ionization equilibrium. This latter case leads to the temperature of the plasma being localized within the values where hydrogen (circle D) and helium are ionized (circle C).

There are several physical processes that lead to these differences in the upper chromosphere. These processes will be described in detail in a follow up paper. In short, the expanding cold bubbles, which have large ambipolar diffusion, are heated and, it is through these bubbles that the photospheric horizontal magnetic field is diffused into the chromosphere. From there, the field expands into the corona pushing chromospheric material to higher layers. A large percentage of the magnetic stresses due to this expansion of the magnetic field, waves and footpoint motions are dissipated in the cold bubble. Moreover, for GOL-OS, the diffused field lines coming from the photosphere interact and dissipate magnetic energy in the upper chromosphere into thermal energy driving temperatures up to $\sim 10^4$ K. The various regions of the atmosphere will impact how mechanical and magnetic energy is transported to greater heights in different manners, since different processes dominate in the various regions. The partial ionization interaction effects also produce a slightly more dynamic upper-chromosphere, transition region and lower corona as a result of the expansion of the diffused magnetic field and the interaction of this expanding magnetic field lines with the ambient magnetic field. These simulations need to be recalculated taking into account a better ion-neutral collision frequency calculation (Vranjes & Krstic 2013). Including the effects of partial ionization is important to many aspects of the physics of the chromosphere and corona, in particular those related to the evolution of the magnetic field and the energization of the outer atmosphere's plasma.

5. Summary

Including the effects of partial ionization is important to many aspects of the physics of the chromosphere and corona, in particular those related to the evolution of the magnetic field and the energization of the outer atmospheres plasma, as the interaction between ionized particles and neutral particles has important consequences on the magneto-thermodynamics. In particular, our simplified 2D radiative MHD models reveal that the ambipolar diffusion dissipates magnetic energy and increases the minimum temperature in the chromosphere and the upper chromosphere may get heated and expanded over a greater range of heights.

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